

Upsilon Decays into Scalar Dark Matter

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We examine decays of a spin-1 bottomonium into a pair of light scalar Dark Matter (DM) particles, assuming that Dark Matter is produced due to exchange of heavy degrees of freedom. We perform a model-independent analysis and derive formulae for the branching ratios of these decays. We confront our calculation results with the experimental data. We show that Dark Matter production in Υ decays may lead to constraints on parameters of the models containing a light spin-0 DM particle.

1. Introduction

We consider the possibility of using of Υ meson decays with missing energy, to test the models with a light spin-0 DM particle. Dark Matter search in heavy meson (and in particular Υ meson) decays may be complimentary to such experiments as DAMA [1, 2, 3], CDMS [4] and XENON [5, 6], which rely on the measurement of kinematic recoil of nuclei in DM interactions and lose (for cold DM particles) sensitivity with decreasing mass of the WIMP, as the recoil energy becomes small.

So far, Υ meson decays into Dark Matter have been considered within the models, where DM particles interaction with an ordinary matter is mediated by some light degree of freedom [7, 8, 9]. Apart from desire of having DM annihilation enhancement (due to a light intermediate resonance) and thus having no tension with the DM relic abundance condition [10]-[13], it is also known that Υ meson SM decay is predominantly due to strong interactions. Thus, the WIMP production branching ratio, in general, is greatly suppressed compared to relevant weak B decays, and in particular to $B \rightarrow K + \text{invisible}$ transition [14, 15]. In light of this, it might seem natural to concentrate only on the models within which Dark Matter production in Υ decays is enhanced due to exchange of a light particle propagator.

Yet, our aim is to study $\Upsilon(1S)$ decay into a pair of spin-0 DM particles, $\Upsilon(1S) \rightarrow \Phi\Phi^*$, and $\Upsilon(3S)$ decay into a pair of spin-0 DM particles and a photon, $\Upsilon(3S) \rightarrow \Phi\Phi^*\gamma$, within the models where light Dark Matter interaction with an ordinary matter is due to exchange of *heavy* particles (with masses exceeding the bottomonium mass). These models may be free of tension related to satisfying the DM relic abundance constraint as well [10, 14, 16, 17, 18]. Also, new experimental data on Υ decays into invisible states have been reported by the BaBar collaboration [19, 20]. According to these data,

$$B(\Upsilon(1S) \rightarrow \text{invisible}) < 3 \times 10^{-4} \quad (1.1)$$

and

$$B(\Upsilon(3S) \rightarrow \gamma + \text{invisible}) < (0.7 - 31) \times 10^{-6} \quad (1.2)$$

where the interval in the r.h.s. of eq. (1.2) is related to the choice of the final state missing mass. These bounds are significantly stronger than those on invisible $\Upsilon(1S)$ decays (with or without a photon emission), reported previously by Belle and CLEO [21, 22] and quoted by Particle Data Group [23]. We show that BaBar experimental data on Υ meson invisible decays may constrain the parameter space of light scalar Dark Matter models, even if there is no Dark Matter production enhancement due to light intermediate states.

We also illustrate that the study of Dark Matter production in Υ decays allows us to test regions of parameter space of light spin-0 DM models that are inaccessible for B meson decays with missing energy. It is also worth mentioning that Υ decays are sensitive to a wider range of WIMP mass than B decays. Thus, the study of WIMP production in Υ decays is complementary to that for B meson decays.

At the energy scales, associated with Υ decays, heavy intermediate degrees of freedom may be integrated out, thus leading to a low-energy effective theory of four-particle interactions. Our strategy would be deriving first model-independent formulae for the $\Upsilon(1S) \rightarrow \Phi\Phi^*$ and $\Upsilon(3S) \rightarrow \Phi\Phi^*\gamma$ branching ratios within the low-energy effective theory. Then, we confront our predictions with the experimental data, deriving model-independent bounds in terms of the Wilson operator expansion coefficients, as the parameters that carry the information on an underlying New Physics model. Finally, within a given model, using the matching conditions for the Wilson coefficients, we translate these bounds into those on the relevant parameters of the considered model.

The talk is based on the results presented in [24] and [25].

2. Model-Independent Analysis

We treat Υ states - neglecting the sea quark and gluon distributions - as bound states of $b\bar{b}$ valence quark-antiquark pair that annihilates - with or without emission of a photon - into a pair of Dark Matter particles. To this approximation, the relevant low-

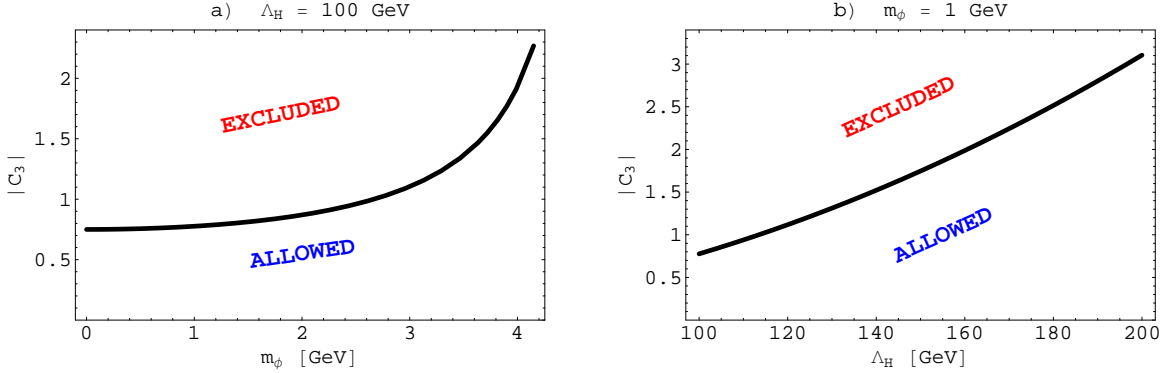


Figure 1: Upper bound on $|C_3|$ a) as a function of m_Φ , for $\Lambda_H = 100 \text{ GeV}$, b) as a function of Λ_H , for $m_\Phi = 1 \text{ GeV}$.

energy effective Hamiltonian may be written as

$$H_{eff} = \frac{2}{\Lambda_H^2} \sum_i C_i O_i \quad (2.1)$$

where Λ_H is the heavy mass and

$$\begin{aligned} O_1 &= m_b (\bar{b} b) (\Phi^* \Phi), \quad O_2 = im_b (\bar{b} \gamma_5 b) (\Phi^* \Phi), \\ O_3 &= (\bar{b} \gamma^\mu b) (\Phi^* i \vec{\partial}_\mu \Phi), \quad O_4 = (\bar{b} \gamma^\mu \gamma_5 b) (\Phi^* i \vec{\partial}_\mu \Phi) \end{aligned} \quad (2.2)$$

with $\vec{\partial} = 1/2(\vec{\partial} - \overleftarrow{\partial})$. It is worth noting that with the notations used in (2.1) and (2.2), all the operators O_i , $i=1, \dots, 4$, are Hermitean, thus all the Wilson coefficients C_i must be real. If DM consists of particles that are their own antiparticles, then only first two operators in eq. (2.2) would contribute.

For DM field being a complex scalar state, $\Upsilon(1S) \rightarrow \Phi\Phi^*$ branching ratio is given by the following expression:

$$\begin{aligned} B(\Upsilon(1S) \rightarrow \Phi\Phi^*) &= \frac{\Gamma(\Upsilon(1S) \rightarrow \Phi\Phi^*)}{\Gamma_{\Upsilon(1S)}} = \\ &= \frac{C_3^2}{\Lambda_H^4} \frac{f_{\Upsilon(1S)}^2}{48\pi\Gamma_{\Upsilon(1S)}} \left[M_{\Upsilon(1S)}^2 - 4m_\Phi^2 \right]^{3/2} \end{aligned} \quad (2.3)$$

where m_Φ is the DM particle mass and $\Gamma_{\Upsilon(1S)}$, $M_{\Upsilon(1S)}$, $f_{\Upsilon(1S)}$ are the $\Upsilon(1S)$ total width, mass and decay constant respectively. Only operator O_3 contributes to the decay rate.

For Φ being a self-conjugate spin-0 state, $\Phi = \Phi^*$, contribution of O_3 vanishes as well, as noted above. Thus, one has

$$B(\Upsilon(1S) \rightarrow \Phi\Phi) = 0 \quad (2.4)$$

This result is related to the fact that the final DM particle pair state must be a P-wave, which is impossible due to the Bose-Einstein symmetry of identical spin-0 particles. In what follows, $\Gamma(\Upsilon(1S) \rightarrow \Phi\Phi)$ must also vanish in higher orders in $1/m_b$ operator product expansion.

Thus, provided that DM pair production is the dominant invisible channel (the neutrino background may be neglected [19, 21, 24]), the signal for $\Upsilon(1S) \rightarrow \text{invisible}$ decay would imply that the light spin-0 DM field has a complex nature. No evidence for the $\Upsilon(1S) \rightarrow \text{invisible}$ mode may lead to some constraints on the parameters of the models with light complex scalar Dark Matter.

Yet, in order to derive such constraints, the experimental limit on the $\Upsilon(1S) \rightarrow \text{invisible}$ mode must be strong enough. Indeed, using the numerical values of $\Gamma_{\Upsilon(1S)}$, $M_{\Upsilon(1S)}$ and $f_{\Upsilon(1S)}$ [23], [26], one may rewrite eq. (2.3) as

$$\begin{aligned} B(\Upsilon(1S) \rightarrow \Phi\Phi^*) &\approx 5.3 \times 10^{-4} C_3^2 \times \\ &\times \left(\frac{100 \text{ GeV}}{\Lambda_H} \right)^4 \left[1 - \frac{4m_\Phi^2}{M_{\Upsilon(1S)}^2} \right]^{3/2} \end{aligned} \quad (2.5)$$

In what follows, the relevant experiments must be sensitive (at least) to $B(\Upsilon(1S) \rightarrow \text{invisible}) \sim 10^{-4}$.

This sensitivity has been reached by the BaBar experiment [19], as it follows from the bound on $B(\Upsilon(1S) \rightarrow \text{invisible})$, given by eq. (1.1). Substituting (1.1) into (2.5), one derives the following constraint on $|C_3|$ as a function of m_Φ and Λ_H :

$$|C_3| < 0.75 \left(\frac{\Lambda_H}{100 \text{ GeV}} \right)^2 \left(1 - \frac{4m_\Phi^2}{M_{\Upsilon(1S)}^2} \right)^{-3/4} \quad (2.6)$$

The behavior of the upper bound on $|C_3|$ with the DM particle mass and the heavy mass is presented in Fig.'s 1a) and 1b) respectively. For $m_\Phi < 3 \text{ GeV}$ and $\Lambda_H \simeq 100 \text{ GeV}$, this bound may be translated into constraints on the relevant couplings of models with a light complex spin-0 DM field. For $m_\Phi > 3 \text{ GeV}$ or for $\Lambda_H > 100 \text{ GeV}$, further improvement of the experimental sensitivity to $\Upsilon(1S) \rightarrow \text{invisible}$ transition is necessary.

Note that even for $m_\Phi < 3 \text{ GeV}$ and $\Lambda_H \simeq 100 \text{ GeV}$, bound (2.6) still allows C_3 to be of the

order of unity. It seems to be very unlikely to saturate such a (rather weak) bound, if within a given New Physics model, $\Upsilon(1S) \rightarrow \Phi\Phi^*$ transition is loop-induced. Thus, bound (2.6) on $|C_3|$ seems to be useful only within the models, where $\Upsilon(1S) \rightarrow \Phi\Phi^*$ decay occurs at tree level.

Further discussion of the scenarios with complex scalar Dark Matter goes beyond the scope of this talk. Yet, in ref. [24] we use bound (2.6) to derive constraints on the parameter space of the models with mirror fermions, in particular on that of the MSSM with gauge mediated SUSY breaking and the DM particle in the hidden sector, while mirror fermions being connectors between the hidden and the MSSM sectors. We derive for the first time bounds on the couplings of b-quark interactions with its mirror counterpart, as functions of the WIMP mass and the mirror fermion mass.

For the scenarios with the DM particle being its own antiparticle, the decay $\Upsilon(3S) \rightarrow \Phi\Phi\gamma$ is relevant. As the neutrino background is negligible [24], $\Upsilon(3S) \rightarrow \Phi\Phi\gamma$ may be the dominant channel in the $\Upsilon(3S) \rightarrow \gamma + \text{invisible}$ mode.

As it was mentioned above, the strongest experimental constraint on $B(\Upsilon \rightarrow \gamma + \text{invisible})$ is that derived by the BaBar collaboration [20] and given by eq. (1.2). It may also be written as

$$B(\Upsilon(3S) \rightarrow \gamma + \text{invisible}) < 3 \times 10^{-6} \quad (2.7)$$

for the final state missing mass being less than 7 GeV, or approximately $\sqrt{s} \lesssim M_{\Upsilon(3S)}/\sqrt{2}$. We have to note however that this bound has been derived assuming that $\Upsilon(3S) \rightarrow \gamma + \text{invisible}$ transition is mediated by an intermediate resonant Higgs state A^0 that exists in the MSSM extensions with an additional Higgs singlet [7, 13, 27]. For the models with non-resonant DM production considered here, bound (2.7) may be used only to get a *preliminary estimate* of possible constraints on parameters of light spin-0 self-conjugate DM models¹. More rigourously, the experimental analysis performed in ref. [20] should be extended to the cases, when the emitted photon energy is non-monochromatic and is in the range $0 < \omega < M_{\Upsilon(3S)}/2 - 2m_\Phi^2/M_{\Upsilon(3S)}$. To our knowledge, this work is in progress now².

Within self-conjugate scalar DM scenarios, the partially integrated branching ratio for $\Upsilon(3S) \rightarrow \Phi\Phi\gamma$

¹As for the earlier bounds on $B(\Upsilon \rightarrow \gamma + \text{invisible})$ [23], the only existing constraint for the case of two-particle invisible states and the photon having non-monochromatic energy, $B(\Upsilon(1S) \rightarrow \gamma + X\bar{X}) < 10^{-3}$ [22], is too weak to have any use. Due to the factor $\alpha/(4\pi) \approx 5.8 \times 10^{-4}$, $B(\Upsilon \rightarrow \Phi\Phi\gamma)$ is always below this limit.

²Author is grateful to Yu. Kolomensky for this information.

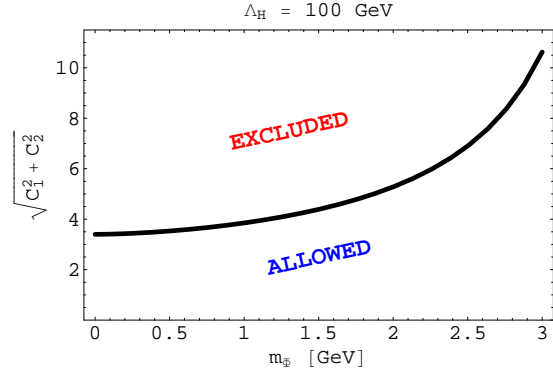


Figure 2: Upper bound on $\sqrt{C_1^2 + C_2^2}$ as a function of m_Φ , for $\Lambda_H = 100 \text{ GeV}$.

decay is given by

$$B(\Upsilon(3S) \rightarrow \Phi\Phi\gamma)_{|s < s_{max}} = \frac{(C_1^2 + C_2^2)}{\Lambda_H^4} \frac{\alpha}{4\pi} \frac{f_{\Upsilon(3S)}^2}{54\pi\Gamma_{\Upsilon(3S)}M_{\Upsilon(3S)}} \times \\ \times \left[\left(2M_{\Upsilon(3S)}^2 - s_{max} + 2m_\Phi^2 \right) \sqrt{s_{max}(s_{max} - 4m_\Phi^2)} - \right. \\ \left. - 8m_\Phi^2 (M_{\Upsilon(3S)}^2 - m_\Phi^2) \ln \left(\frac{\sqrt{s_{max}} + \sqrt{s_{max} - 4m_\Phi^2}}{2m_\Phi} \right) \right] \quad (2.8)$$

For $s_{max} = M_{\Upsilon(3S)}^2/2$, using the numerical values of the $\Upsilon(3S)$ mass, total width and decay constant [23, 24], one may rewrite eq. (2.8) in the following form:

$$B(\Upsilon(3S) \rightarrow \Phi\Phi\gamma)_{|s < M_{\Upsilon(3S)}^2/2} = \\ = 2.6 \times 10^{-7} (C_1^2 + C_2^2) \left(\frac{100 \text{ GeV}}{\Lambda_H} \right)^4 f(x_\Phi) \quad (2.9)$$

where $x_\Phi = m_\Phi^2/M_{\Upsilon(3S)}^2$ and

$$f(x_\Phi) = \left(1 + \frac{4}{3}x_\Phi \right) \sqrt{1 - 8x_\Phi} - \\ - \frac{32}{3}x_\Phi(1 - x_\Phi) \ln \left(\frac{1 + \sqrt{1 - 8x_\Phi}}{2\sqrt{2}\sqrt{x_\Phi}} \right) \quad (2.10)$$

At first glance, it may seem that $\Upsilon(3S) \rightarrow \Phi\Phi\gamma$ branching ratio is far out of reach of the BaBar experimental sensitivity, for a reasonable choice of C_1 and C_2 . Notice, however, that within certain models with light spin-0 Dark Matter, the Wilson coefficients C_1 and/or C_2 may be enormously large, as they contain some enhancement factors, such as the ratio $\Lambda_H/m_b \gg 1$ - due to the mass term in the numerator of a heavy fermion propagator, or the Higgs vev's ratio $\tan\beta$ (with the latter being, say, $\sim m_t/m_b \gg 1$) - due to DM particle pair production via exchange of a heavy non-SM Higgs degree of freedom.

These enhancement factors can make C_1 and/or C_2 to be $\gtrsim 10$ and hence $B(\Upsilon(3S) \rightarrow \Phi\Phi\gamma)$ to be \sim

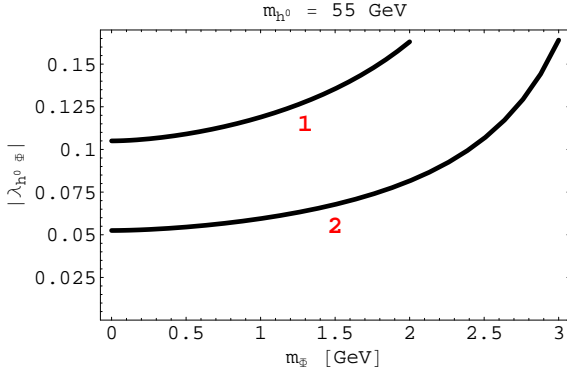


Figure 3: Upper bound on $|\lambda_{h^0 \Phi}|$ as a function of m_Φ , for $m_{h^0} = 55 \text{ GeV}$ and $\tan \beta = 20$ (line 1), $\tan \beta = 40$ (line 2).

$10^{-5} - 10^{-4}$, i.e. significantly exceeding bound (2.7) on $B(\Upsilon \rightarrow \gamma + \text{invisible})$. Using bound (2.7) yields

$$\sqrt{C_1^2 + C_2^2} < 3.4 \left(\frac{\Lambda_H}{100 \text{ GeV}} \right)^2 f^{-1/2}(x_\Phi) \quad (2.11)$$

The behavior of the bound on $\sqrt{C_1^2 + C_2^2}$ with the DM particle mass is presented in Fig. 2 for $\Lambda_H = 100 \text{ GeV}$.

In the next section we consider the simplest possible model where the Wilson coefficients are made enormously large due to an enhancement factor. We transform the constraint on $\sqrt{C_1^2 + C_2^2}$ into that on the relevant parameters of the model.

3. Dark Matter Model with two Higgs Doublets (2HDM)

In this section we consider the two-Higgs doublet model (2HDM) with a gauge singlet real scalar DM particle. The DM interaction part of Lagrangian, relevant for our analysis, may be written as [14]

$$-\mathcal{L} = \frac{m_0^2}{2} \Phi^2 + \lambda_1 \Phi^2 |H_1|^2 + \lambda_2 \Phi^2 |H_2|^2 + \lambda_3 \Phi^2 (H_1 H_2 + h.c) \quad (3.1)$$

Following refs. [14, 28], we consider type-II version of 2HDM, where H_1 generates masses of down-type quarks and charged leptons, whereas H_2 generates masses of up-type quarks. We assume for the Higgs vev's ratio to be large, i.e. $v_2/v_1 \equiv \tan \beta \gg 1$.

Three coupling constants entering Lagrangian (3.1) are unknown parameters. Constraints on the couplings λ_1 and λ_2 may be derived from the study of $B \rightarrow K + \text{invisible}$ transition. Those, in general, are strong enough: combined with the ones coming from the DM relic abundance condition, they rule out a wide range of the WIMP mass for the scenarios with λ_1 and/or λ_2 dominant [14].

Yet, due to cancellation effects in the relevant diagrams, WIMP pair production rate in B meson decays is insensitive to the value of λ_3 [14]. The scenarios with λ_3 dominant, or at least non-negligible, are thus far unconstrained. Study of DM production in Υ decays enables one to constrain these scenarios, inaccessible by B meson decays with missing energy.

The model contains two CP-even, one CP-odd and two complex charged physical Higgs degrees of freedom. To the leading order in the perturbation theory, $\Upsilon(3S) \rightarrow \Phi \Phi \gamma$ transition occurs at tree level by exchange of a single Higgs boson. To this approximation, only the CP-even Higgs states, h^0 and H^0 , are relevant for our analysis.

One of the CP-even Higgs bosons (not-necessarily the lightest one) is Standard Model-like: its phenomenology is similar to that of the SM Higgs boson and the experimental bound on its mass is close to the SM limit³ (see [23] and references therein). Contribution of the diagrams with the SM-like Higgs boson exchange to the $\Upsilon(3S) \rightarrow \Phi \Phi \gamma$ decay amplitude has no any enhancement factor - thus, we may further disregard the Dark Matter interaction with the SM-like CP-even Higgs.

The other CP-even Higgs boson is "New-Physics (NP) like": its phenomenology differs drastically from that of the SM Higgs boson [31]. If Dark Matter is produced due to exchange of this Higgs particle, $\Upsilon(3S) \rightarrow \Phi \Phi \gamma$ decay amplitude is enhanced by $\tan \beta$ factor. As discussed above, in this case the decay branching ratio is within the reach of the BaBar experimental sensitivity.

If the NP-like Higgs is the lightest one, its mass may be much below the Standard Model experimental limit: according the existing experimental data [32, 33], $m_{h^0} > 55 \text{ GeV}$ or $m_{h^0} < 1 \text{ GeV}$ in the general type II 2HDM. As it was mentioned above, the light Higgs scenario is beyond the scope of the present paper, thus we assume here that $m_{h^0} > 55 \text{ GeV}$. Notice, however, that this bound is derived, provided that no invisible Higgs decay mode exists. On the other hand, if the NP-like Higgs invisible decay mode is dominant, it may escape detection. No bound on m_{h^0} , to our best knowledge, exists in that case.

Within the considered model with light scalar Dark Matter, analysis of the $\Upsilon \rightarrow \Phi \Phi \gamma$ mode may restrict the scenarios with an invisibly decaying lightest Higgs boson by putting severe constraints on the $h^0 \Phi \Phi$ interaction coupling, $\lambda_{h^0 \Phi}$. For $\tan \beta \gg 1$ it is given by

$$\lambda_{h^0 \Phi} \approx \lambda_3 + (\lambda_1 + \lambda_2) \cos \beta \quad (3.2)$$

³Also, if the SM Higgs decays predominantly invisibly, the SM lower experimental bound is distorted by a few GeV only [29, 30].

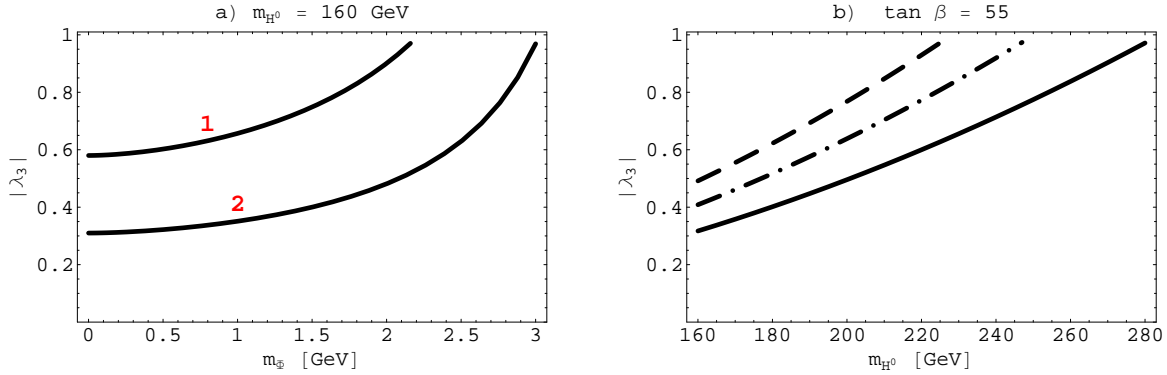


Figure 4: Upper bound on $|\lambda_3|$, a) as a function of m_Φ for $m_{H^0} = 160$ GeV and $\tan \beta = 30$ (line 1), $\tan \beta = 55$ (line 2), b) as a function of m_{H^0} for $\tan \beta = 55$ and $m_\Phi = 100$ MeV (solid line), $m_\Phi = 1.5$ GeV (dashed-dotted line), $m_\Phi = 2$ GeV (dashed line).

The last term in the r.h.s. of (3.2), although being suppressed by a factor of $\cos \beta \approx 1/\tan \beta$, must be retained because of possible hierarchy in the values of λ_3 and λ_1 or λ_2 . Scenarios with such a hierarchy may be of importance: using the matching conditions for the Wilson coefficients,

$$C_1 = -\frac{\lambda_{h^0\Phi}}{2} \tan \beta, \quad C_2 = 0, \quad \Lambda_H = m_{h^0}, \quad (3.3)$$

and bound (2.11) on $\sqrt{C_1^2 + C_2^2}$, one finds that $\lambda_{h^0\Phi}$ is constrained to be $O(1/\tan \beta)$, if h^0 mass approaches to its lower limit, $m_{h^0} = 55$ GeV. We illustrate this in Fig. 3 for $\tan \beta = 20$ and $\tan \beta = 40$.

For $\lambda_{h^0\Phi}$ being so strongly constrained, it seems to be very unlikely that h^0 would escape detection and its mass be below 55 GeV. More rigorously, however, detailed re-analysis of the Higgs production and decay rates, including that of $h^0 \rightarrow \Phi\Phi$, should be performed.

One can show [24] that when choosing $\tan \beta$ sufficiently large, bound on $\lambda_{h^0\Phi}$ may still be rigorous, if the h^0 mass is heavier than 55 GeV. Thus, within the type II 2HDM with a light spin-0 Dark Matter, study of $\Upsilon \rightarrow \Phi\Phi\gamma$ decay channel may lead to severe constraints on the lightest CP-even Higgs invisible decay coupling, if that Higgs is New-Physics like.

The remarkable feature of the model is that constraints on the parameter space are derived even if the NP-like Higgs is the heaviest CP-even one. We illustrate this choosing $m_{H^0} \geq 160$ GeV - the existing theoretical upper bound on the SM-like Higgs boson [34] and exclusion of the SM Higgs mass interval (160 – 170) GeV with 95% C. L. by the CDF and D0 data [35] allow us to infer that above 160 GeV, the CP-even Higgs boson is presumably the heaviest one and NP-like. The matching conditions for the Wilson coefficients are now

$$C_1 = \frac{-\lambda_3 \tan \beta}{2}, \quad C_2 = 0, \quad \Lambda_H = m_{H^0} \quad (3.4)$$

Thus, bound (2.11) on $\sqrt{C_1^2 + C_2^2}$ may be transformed into that on $|\lambda_3|$ as a function of the WIMP mass, $\tan \beta$ and the heaviest CP-even Higgs mass.

As one can see from Fig. 4a), for $m_{H^0} = 160$ GeV and $\tan \beta = 30$, λ_3 is constrained to be of order of the SM weak coupling or smaller ($|\lambda_3| \lesssim 0.65$), if $m_\Phi \lesssim 1$ GeV. Also, for the same choice of the Higgs mass and $\tan \beta$, $|\lambda_3|$ is to be less than one, if the WIMP mass is less than 2 GeV. Bound on $|\lambda_3|$ is significantly more rigorous for higher values of $\tan \beta$. For instance, if choosing $\tan \beta = m_t(m_t)/m_b(m_t) \approx 55$, one gets $|\lambda_3| \lesssim 0.35$ and $|\lambda_3| < 0.5$ for $m_\Phi \lesssim 1$ GeV and $m_\Phi \simeq 2$ GeV respectively.

The restrictions on $|\lambda_3|$ are essential also for higher values of the heaviest CP-even Higgs mass: for $\tan \beta = 55$, they are still of the interest up to $m_{H^0} \simeq 280$ GeV, as one can see from Fig. 4b).

Thus, within the type II 2HDM with a scalar Dark Matter, for large $\tan \beta$ scenario, Υ meson decay into a Dark Matter particles pair and a photon, $\Upsilon \rightarrow \Phi\Phi\gamma$, may be used to derive essential constraints on the parameters of the model, which otherwise cannot be tested by B meson decays with invisible outgoing particles.

4. Conclusions and Summary

Thus, spin-0 Dark Matter production in Υ meson decays has been investigated. We restricted ourselves by consideration of the models where the decays occur due to exchange of heavy non-resonant degrees of freedom.

We performed our calculations within low-energy effective theory, integrating out heavy degrees of freedom. This way we derived model-independent formulae for the considered branching ratios. We used these formulae to confront our theoretical predictions with existing experimental data on invisible Υ decays, both in a model-independent way and within particu-

lar models. It has been shown that within the considered class of models, DM production rate in Υ decays is within the reach of the present experimental sensitivity. Thus, Υ meson decays into Dark Matter, with or without a photon emission, may be used to constrain the models with a GeV or lighter spin-0 DM.

Experimental constraints on the $\Upsilon(3S) \rightarrow \gamma + \textit{invisible}$ mode are derived assuming that Dark Matter is produced by exchange of a light resonant scalar state. Within the scenarios with non-resonant DM production considered here, these constraints may be used only to make preliminary estimates of possible bounds on the parameters of the models. Yet, those estimates show that these bounds may be rigorous enough; besides, they are derived within the least restrained presently light scalar DM scenarios. Our goal is thus to encourage the experimental groups to analyze the experimental data on $\Upsilon \rightarrow \gamma + \textit{invisible}$ also for the case of non-monochromatic photon emission and spin-0 invisible states.

So, from our analysis one may conclude that Dark Matter production in Υ meson decays may serve as an interesting alternative to commonly used DM search methods, capable of providing a valuable information on DM particles, if those turn to have a mass of the order of a few GeV or smaller.

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